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Asymmetric Inflationary Reheating and the Nature of Mirror Universe

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Abstract

The existence of a shadow world (or mirror universe) with matter and forces identical to that of the visible world but interacting with the latter only via gravity can be motivated by superstring theories as well as by recent attempts to understand the nature of a sterile neutrino needed if all known neutrino data are to be consistent with each other. A simple way to reconcile the constraints of big bang nucleosynthesis in such a theory is to postulate that the reheating temperature after inflation in the mirror universe is lower than that in the visible one. We have constructed explicit models that realize this proposal and have shown that the asymmetric reheating can be related to a difference of the electroweak symmetry breaking scales in the two sectors, which is needed for a solution of the neutrino puzzles in this picture. Cosmological implications of the mirror matter are also discussed.

1. Introduction

The $E_8 \times E'_8$ string theories indicate an interesting possibility that microphysics of the early universe has two parallel components with identical matter and force structure which communicate only through gravity [1]. In a recent paper [2], two of us have suggested that the same idea may be motivated by the existing neutrino data (for a different model see also [3]).¹⁾ The reasoning goes as follows: the simplest way to reconcile the present neutrino puzzles is to invoke a fourth neutrino which is sterile with respect to the weak interactions and extremely light (with the mass in the milli-eV range) [5]. The solar neutrino problem (SNP) can be then explained by the MSW mechanism operating between the ν_e and ν_s whereas the atmospheric neutrino data is explained by the oscillation between the nearly degenerate ν_μ and ν_τ states with mass of 2.5 eV which also provide a cosmological hot dark matter (HDM). The recent LSND results can be also explained by small oscillations between ν_e and ν_μ . From a theoretical point of view, the lightness of the sterile neutrino would be easier to understand if it could be subjected to the same kind of symmetry reasoning that keeps the known neutrinos light, i.e. accidental B-L symmetry possibly broken by gravity or a local B-L symmetry broken at some very high scale as in the usual seesaw mechanism. If one postulates a mirror universe with identical gauge and matter content [4], the neutrinos of the mirror universe become subjected to the mirror B-L symmetry and remain light. In particular, the sterile state ν_s can be the mirror partner ν'_e of the usual electron neutrino. In fact it is shown in Ref. [2] that using only one input that the electroweak symmetry breaking scale in the mirror universe is a factor $\zeta \sim 30$ higher than the usual electroweak scale, simultaneously gives both the desired mass and mixing range for the MSW oscillation $\nu_e - \nu'_e$ to be successful in solving the solar neutrino problem. In addition, if ν_μ and ν_τ have masses in eV range constituting thereby HDM, their mirror partners ν'_μ and ν'_τ being factor of $\zeta^2 \sim 1000$ heavier, will have masses in keV range and thus can provide the warm dark matter.

An immediate challenge to this approach is to reconcile the constraints of big bang nucleosynthesis (BBN) on the energy density of the universe at the BBN epoch, which translates into constraints on the number of light neutrinos N_ν [6]. In a theory such as ours, since all mirror neutrinos (and photons) are light, they could apriori give a large contribution considerably exceeding $N_\nu = 3$. Therefore, for our idea to be viable, the contribution of the light mirror particles to the energy density at the BBN epoch must be appropriately reduced. The idea to achieve this goal suggested in Ref. [2] was to postulate an asymmetric postinflationary reheating of the two universes. In particular, if the mirror universe reheats to a lower temperature than our universe, then BBN constraint can be satisfied. The purpose of the present paper is to present realizations of this idea in the context of gauge models and then study cosmological consequences of this hypothesis. We also discuss the state of the mirror universe at the present epoch. In particular, we argue that for the case when the electroweak symmetry breaking in the mirror sector is about 30 times larger than that of the visible universe, it is likely that the mirror baryonic matter would consist entirely of the mirror hydrogen which might be the only stable mirror atom. Its implications for the microlensing experiments are also discussed.

¹⁾The concept of a hidden mirror world has been considered also in several earlier papers [4].

2. Mirror world: a brief review and cosmological implications

Having in mind the $E_8 \times E'_8$ string model, one can imagine that below the Planck (string) scale field theory is given by a product of two identical gauge groups $G \times G'$ with identical particle contents, and there is a discrete symmetry $P(G \leftrightarrow G')$ interchanging all particles in corresponding representations of G (which we consider as visible world) and G' (which we call mirror world). It guarantees that all bare coupling constants (gauge, Yukawa, Higgs) have the same values in both sectors. We also assume that the two worlds communicate only through gravity and possibly also via a heavy gauge singlet matter field. At some scale below M_{Pl} the gauge symmetry breaks down to $G_{SM} \times G'_{SM}$, where $G_{SM} = SU(3) \times SU(2) \times U(1)$ stands for the standard model incorporating quarks and leptons $q_i, u_i^c, d_i^c; l_i, e_i^c$ and Higgs doublet ϕ , and $G'_{SM} = [SU(3) \times SU(2) \times U(1)]'$ is its mirror counterpart with analogous particle content: $q'_i, u_i'^c, d_i'^c; l'_i, e_i'^c$ and ϕ' ($i = 1, 2, 3$ is a family index).²⁾ P parity remains unbroken at this stage, so that the coupling constants of the two sectors evolve down in energy from common values. Let us also assume that there exists a mechanism that spontaneously breaks P parity at lower energies and thus allows the two electroweak scales $\langle \phi' \rangle = v'$ and $\langle \phi \rangle = v$ to be different: we assume that $v' \gg v$. As far as the Yukawa couplings have the same values in both systems, the mass and mixing pattern of the charged fermions in the mirror world is completely analogous to that of the visible one, but with all fermion masses scaled up by the factor $\zeta = v'/v$. The masses of gauge bosons and higgses are also scaled as $M_{W', Z', \phi'} = \zeta M_{W, Z, \phi}$ while photons and gluons remain massless in both sectors.

As for the neutrino masses, they can emerge only via operators bilinear in the Higgs fields and cutoff by a large scale M , which can be effectively induced for example via the seesaw mechanism. On general grounds, by assuming that P parity breaks at lower energies ($E \ll M$), these operators can be written as

$$\frac{h_{ij}}{M}(l_i\phi)C(l_j\phi) + \frac{h'_{ij}}{M}(l'_i\phi')C(l'_j\phi') + \text{h.c.} \quad (1)$$

where C is a charge conjugation matrix. In order to deal with the present neutrino data, one can assume further that $M \sim 10^{13}$ GeV and that the $O(1)$ coupling constants h_{ij} obey an approximate $L_e + L_\mu - L_\tau$ symmetry [2]. Thus ν_e and ν'_e are left massless, $\nu_{\mu, \tau}$ get almost degenerate masses $m \sim v^2/M \sim \text{few eV}$, and the masses of their mirror partners $\nu'_{\mu, \tau}$ are $m' = \zeta^2 m$. The ν_e and ν'_e states then can get masses through the gravity induced Planck scale effects [7, 8] which explicitly violate the global lepton number, and also mix the neutrino states of two sectors [8]. The relevant operators are:

$$\frac{\alpha_{ij}}{M_{Pl}}(l_i\phi)C(l_j\phi) + \frac{\alpha'_{ij}}{M_{Pl}}(l'_i\phi')C(l'_j\phi') + \frac{\beta_{ij}}{M_{Pl}}(l_i\phi)C(l'_j\phi') + \text{h.c.} \quad (2)$$

²⁾It is essential that at higher energies both $SU(3) \times SU(2) \times U(1)$ gauge factors are embedded into simple gauge groups. Otherwise kinetic terms of the two $U(1)$ gauge fields could mix and this would impart arbitrary electric charges to the particles [4]. In the spirit of this proposal, one may therefore envisage that the gauge groups G and G' are identical and simple (e.g. $SU(5)$, $SU(6)$, $SO(10)$ or any other GUT subgroup of E_8). Even in this case, the kinetic mixing between usual and mirror photons would arise from radiative effects in the presence of the mixed representations of $G \times G'$. Bearing in mind the possible $E_8 \times E'_8$ string origin for our models, here we exclude such mixed representations.

with constants $\alpha, \beta \sim 1$. Then ν_e and ν'_e acquire masses respectively $\sim \mathring{m}$ and $\sim \zeta^2 \mathring{m}$, while their mixing term is $\sim \zeta \mathring{m}$, where $\mathring{m} = v^2/M_{Pl} = 2.5 \cdot 10^{-6}$ eV. Hence, the $\nu_e - \nu'_e$ oscillation emerges with parameters in the range

$$\delta m^2 \sim (\zeta/30)^4 \times 6 \cdot 10^{-6} \text{ eV}^2, \quad \sin^2 2\theta \sim (30/\zeta)^2 \times 5 \cdot 10^{-3} \quad (3)$$

which for $\zeta \sim 30$ perfectly fit the “small mixing angle” MSW solution to the SNP [9]. More generally, by taking into account the solar model uncertainties, as well as possible order of magnitude spread in the constants α, β , the relevant range for ζ can be extended to $\zeta \sim 10-100$ [2]. Alternatively, for $\zeta \sim 1$ we get $\delta m^2 \sim 10^{-11} \text{ eV}^2$ and $\sin^2 2\theta \sim 1$, which corresponds to another solution for SNP known as ‘just-so’ scenario [10]. Independently of the value of ζ , the values of δm^2 and $\sin^2 2\theta$ given by eq. (3) are safely below the BBN bounds on the active to sterile neutrino oscillation $\nu_e - \nu'_e$ [11] even if a very strong upper bound $\Delta N_\nu < 0.1$ is taken. The same is true for the oscillations $\nu_{\mu,\tau} - \nu'_{\mu,\tau}$, with mixing between them induced by the Planck scale operators (2). However, the oscillations between $\nu_{\mu,\tau}$ and ν'_e with $\delta m^2 \simeq -m^2$ and $\sin^2 2\theta \simeq 4\zeta^2(\mathring{m}/m)^2$ may possess a resonant behaviour in the cosmic plasma, and in accordance with the estimates of [11] we roughly get an upper limit $\zeta < 10^3$. This excludes very high values of the mirror electroweak scale and thus supports our proposal that the P parity breaking is a lower energy phenomenon.

With regard to the two chromodynamics, a big difference between the electroweak scales v' and v will not cause the similar big difference between the confinement scales in two worlds. Indeed, if P parity is valid at higher (GUT) scales, the strong coupling constants in both sectors would evolve down in energy with same values until the energy reaches the value of the mirror-top (t') mass. Below it α'_s will have a different slope than α_s . It is then very easy to calculate the value of the scale Λ' at which α'_s becomes large. This value of course depends on the ratio $\zeta = v'/v$. Taking $\Lambda = 200$ MeV for the ordinary QCD, we find $\Lambda' \simeq 280$ MeV if $\zeta \sim 30$. On the other hand, we have $m'_{u,d} = \zeta m_{u,d} \sim m_s$ so that masses of the mirror light quarks u' and d' do not exceed Λ' . So the condensates $\langle \bar{q}'q' \rangle$ should be formed with approximately the same magnitudes as the usual quark condensates $\langle \bar{q}q \rangle$. As a result, mirror pions should have mass $m'_\pi \simeq \sqrt{m_{u',d'} \langle \bar{q}'q' \rangle}$ comparable to the mass of normal kaons $m_K \simeq \sqrt{m_s \langle \bar{q}q \rangle}$.

As for the mirror nucleons, their masses are approximately 1.5 times larger than that of the usual nucleons. Since $(m'_d - m'_u) \approx 30(m_d - m_u)$ we expect the mirror neutron n' to be heavier than the mirror proton p' by about 150 MeV or so, while the mirror electron mass is $m'_e = \zeta m_e \sim 15$ MeV. Clearly, such a large mass difference cannot be compensated by the nuclear binding energy and hence even bound neutrons will be unstable against β decay $n' \rightarrow p'e'\bar{\nu}'_e$. Thus in the mirror world hydrogen will be the only stable nucleus.

Concerning thermodynamics of the two worlds in the Early Universe, we assume that already at the postinflationary “reheating” stage they are decoupled from each other. As we discuss in next section, once P -invariance is spontaneously broken, it can be violated also in the inflaton couplings to matter. Then the inflaton should decay into visible and mirror particles with different rates, so that after inflation the temperatures of the ordinary (T_R) and mirror (T'_R) thermal bathes would be different (the idea of using inflation to provide a temperature difference between ordinary matter and mirror or other forms of hidden matter was first discussed in ref. [12]). In this way, the present cosmological

abundance of light mirror particles can be suppressed as compared to that of their visible partners.

In the standard cosmology the effective number of the light degrees of freedom at the BBN era is $g_* = 10.75$ as is contributed by photons γ , e^+e^- pairs, and three neutrino species $\nu_{e,\mu,\tau}$, in a good agreement with the observed light element abundances [6]. In the presence of the mirror universe mirror photons γ' and neutrinos $\nu'_{e,\mu,\tau}$ would also contribute the effective number of neutrino species N_ν :

$$\Delta g_* = 1.75 \Delta N_\nu = (2 + 5.25x^4) \left(\frac{T'}{T} \right)^4, \quad x = T'_\nu/T' \quad (4)$$

Here T , T' and T'_ν are respectively the temperatures of γ , γ' and ν' at the BBN era. The value of x is determined by the temperature T'_D at which ν' s decouple from the mirror electromagnetic plasma. T'_D can be expressed through the decoupling temperature of the ordinary neutrinos ($T_{D(\nu_e)} \simeq 2$ MeV and $T_{D(\nu_\mu)} \simeq 3$ MeV). It is scaled as $T'_D \sim \zeta^{4/3}(g_*T^4/10.75T'^4 + g'_*/10.75)^{1/6}T_D$, where the first factor is related to the larger masses of the mirror intermediate bosons and the second one comes because the universe expansion is dominated by the ordinary particles. If $\zeta \simeq 30$ and $T > 2T'$, then mirror neutrinos decouple before the mirror QCD phase transition: $T'_D > \Lambda' \simeq 280$ MeV $\geq m'_{u,d} = \zeta m_{u,d}$, so that the light quarks u' , d' and mirror gluons also contribute along with the electron e' to the heating of γ' . This leads to $x = (4/85)^{1/3}$ and by taking $\Delta N_\nu < 0.1$ we obtain³⁾

$$\frac{T'}{T} \approx 0.95 (\Delta N_\nu)^{1/4} < 0.52 \quad (5)$$

According to this bound the present day abundance of the mirror neutrinos relative to the usual ones $r = n_{\nu'}/n_\nu = (xT'/T)^3$ should be less than 10^{-2} . The usual and mirror neutrinos contribute to the present cosmological density as

$$\sum m_\nu = \Omega_\nu \cdot 94h^2 \text{ eV}, \quad \sum m'_\nu = \zeta^2 \sum m_\nu = r^{-1} \Omega'_\nu \cdot 94h^2 \text{ eV} \quad (6)$$

where h is the Hubble constant in units $100 \text{ Km s}^{-1} \text{ Mpc}^{-1}$. One can assume further that ordinary neutrinos have mass in the eV range and thus form the HDM component with $\Omega_\nu \simeq 0.2$ (as in the model [2] with almost degenerate ν_μ and ν_τ having masses of about 2.5 eV). Then the mirror neutrino masses being factor of ζ^2 larger emerge in the keV range and thus could constitute the warm dark matter (WDM) of the universe. From (6) we get $r = \zeta^{-2}(\Omega'_\nu/\Omega_\nu)$. Then taking a rather conservative bound $\Omega'_\nu < 0.8$ (bearing in mind that other particles like LSP or mirror baryons could also contribute to the present energy density), we obtain the upper bound comparable to that of Eq. (5):

$$\frac{T'}{T} < \frac{1.6}{x\zeta^{2/3}} \approx 0.4 \left(\frac{30}{\zeta} \right)^{\frac{2}{3}} \quad (7)$$

The obtained limits on T'/T can be immediately translated to the limit on the ratio of the postinflationary reheating temperatures. Indeed, the two worlds are decoupled

³⁾One can easily check that this value of x remains constant up to $\zeta \sim 10^5$, and then decreases step-by-step due to contributions of the heavier states μ' , s' etc. On the contrary, for $\zeta \leq 10$ the decoupling of ν' occurs below Λ' and we arrive to the ‘standard’ result $x = T'_\nu/T' = (4/11)^{1/3}$.

and if during the universe expansion both of them evolve adiabatically with separately conserved entropy then we arrive to the nucleosynthesis epoch having ⁴⁾

$$\frac{T'_R}{T_R} = \left(\frac{2 + 5.25x^3}{10.75} \right)^{\frac{1}{3}} \frac{T'}{T} \approx 0.6 \frac{T'}{T} \quad (8)$$

Hence, in the most interesting case the electroweak scale v' in the mirror sector should be by factor $\zeta \sim 30$ larger than the standard electroweak scale $v = 174$ GeV, while the reheating temperature of the mirror universe should be 4-5 times smaller than that of the visible one. In this case cosmological dark matter can consist dominantly of neutrinos. In particular, if the ordinary neutrinos $\nu_{\mu,\tau}$ with the mass of few eV's form the HDM component with $\Omega_\nu \simeq 0.2$, then their mirror partners $\nu'_{\mu,\tau}$ being $\zeta^2 \sim 1000$ times heavier emerge in keV range and with the present abundance smaller by two orders of magnitude than that of normal neutrinos, would constitute the WDM with $\Omega'_\nu \simeq 0.7$. Clearly, their masses satisfy the Tremaine-Gunn limit [13] and thus could constitute dark matter even in dwarf spheroidal galaxies where this limit is most stringent ($m'_\nu > 0.3 - 0.5$ keV). The implications of such mixed HDM+WDM scenario for the large scale structure are rather similar [14] to that of the currently popular HDM+CDM scheme [15] with the cold dark matter (CDM) consisting of heavy ($m \sim 100$ GeV) particles or axionic condensate. However, more detailed observational data on the large scale structure of the universe may make it possible to discriminate between warm and cold dark matter. Moreover, dark matter consisting of sterile neutrinos invalidates direct searches of the CDM candidates via superconducting detectors or axion haloscopes. High energy neutrino fluxes from the Sun and from the Galactic center which are expected from the annihilation of LSP's if they dominate dark matter in the universe, will also be absent. One should however keep in mind that in supersymmetric versions of our scheme CDM as well could exist in the form of the LSP.

An interesting question is what is the amount and form of the mirror baryonic dark matter in the universe. Most likely, baryogenesis in the mirror universe proceeded through the same mechanism as in the visible one and we may expect that the baryon asymmetries (BA) in both worlds should be nearly the same. Since mirror nucleons are not much heavier than the usual ones their fraction in the present energy density, $\Omega_{B'}$, would be about the same as Ω_B , that is around a few percent.

Let us discuss now cosmological evolution of mirror baryonic matter. Since the binding energy of the mirror hydrogen atom is thirty times larger than that of the ordinary hydrogen, its recombination occurs much earlier than the usual recombination era. Hence, the evolution of density fluctuations in the mirror matter would be more efficient than in the visible one. (From the viewpoint of the visible observer mirror baryons behave as a dissipative dark matter.) As a result, one can expect that the distribution of mirror baryons in galactic discs should be more clumped towards the center. It is noteworthy that mirror dark matter may show antibiasing behaviour ($b < 1$) which is considered unphysical for normal dark matter. Since mirror hydrogen is the only stable nucleus

⁴⁾We also assume that initially $g_* = g'_*$ despite different T_R and T'_R , which is natural if $T_R, T'_R \gg v'$. In addition, the relation (8) holds if there are no first order phase transitions. In the presence of the latter the situation could be very much different (see e.g. model of Section 4).

in the mirror world, nuclear burning could not be ignited and luminous (in terms of γ') mirror stars cannot be formed. Therefore, nothing can prevent the sufficiently big protostars from gravitational collapse and in dense regions like galactic cores a noticeable fraction of mirror baryons would collapse into black holes. Recent observational data indeed suggest a presence of giant black holes with masses $\sim 10^{6-7} M_\odot$ in galactic centers. In addition, easier formation of mirror black holes may explain the early origin of quasars.

The remaining fraction of the mirror baryons could fragment into smaller objects like white dwarves (or possibly neutron stars) which can maintain stability due to the pressure of degenerate fermions. For the mirror stars consisting entirely of hydrogen, the Chandrasekhar limit is $M'_{\text{Ch}} \simeq 5.75(m_p/m'_p)^2 M_\odot$, where m_p and $m'_p \simeq 1.5m_p$ are respectively the masses of usual and mirror proton. For smaller mirror objects the evaporation limit should be 2 – 3 orders of magnitude smaller than for the visible ones because the Bohr radius of the mirror hydrogen is 30 times smaller than that of the usual one.

These mirror objects, being dark for the normal observer, could be observed as Machos in the gravitational microlensing experiments (for a review, see e.g. [16]). In principle they can be distinguished from the Machos of the visible world. The latter presumably consist of the dim compact objects (brown dwarves) too light to burn hydrogen, with masses ranging from the evaporation limit ($\sim 10^{-7} M_\odot$) to the ignition limit ($\sim 10^{-1} M_\odot$) [17]. The mass spectrum of mirror Machos extends from the evaporation limit $\sim 10^{-9} M_\odot$ up to the Chandrasekhar limit $M'_{\text{Ch}} \simeq 3 M_\odot$. The present data on the microlensing events are too poor to allow any conclusion on the presence of such heavy (or light) objects. An unambiguous determination of the Macho mass for each event is impossible, and only the most probable mass can be obtained, depending of the spatial and velocity distribution of Machos. The optical depth or the fraction of the sky covered by the Einstein disks of Machos, is nearly independent of their mass: the Einstein disk surface is proportional to M , while the number of deflectors for a given total mass decreases as M^{-1} . However, larger event statistics will allow to find the Macho mass distribution with a better precision.

As noted earlier, the distribution of mirror baryonic matter in galaxies should be more shifted towards their centers as compared to the visible matter. Thus one can expect that mirror stars in our Galaxy would significantly contribute to the microlensing events towards the galactic bulge, while their contribution to the microlensing in halo should be smaller than that of usual Machos. Interestingly, the event rates in the galactic bulge observed by OGLE and MACHO experiments are about twice larger than the expected value deduced from the low mass star population in the Galactic disk [18]. Barring accidental conspiracies like a presence of bar (elongated dense stellar distribution along the line of sight), this can be explained by the contribution of mirror stars, which could naturally increase the optical depth towards the galaxy bulge by factor 2 or so.

In dwarf galaxies mirror baryons may be less concentrated than in spiral ones and may form relatively extended halos. Recent observational data indeed suggest that the distribution of dark matter in dwarves support the idea of dissipative dark matter [19].

3. Asymmetric reheating and asymmetry of electroweak scales

In view of the analysis of the previous section the basic requirements to the model

are to provide a ground state in which both electroweak VEVs are nonzero and different, to provide different inflationary reheating in both worlds, and to suppress all possible interactions which could bring two worlds into thermal equilibrium with each other.

Let us consider a toy model involving two scalars ϕ, ϕ' , analogues to our matter fields. Their Lagrangian, invariant under the discrete transformation $P: \phi \leftrightarrow \phi'$, has the form

$$\mathcal{V}_0(\phi, \phi') = (m_0^2 \phi^2 + h_0 \phi^4) + (m_0'^2 \phi'^2 + h_0' \phi'^4) + a_0 \phi^2 \phi'^2 \quad (9)$$

The simplest way to spontaneously break P -invariance is to introduce a real P -odd scalar η ($P: \eta \rightarrow -\eta$), with nonzero VEV [20]. For our model it is natural to assume that η also plays the role of inflaton. Let us present the potential of η in a generic form $\mathcal{V}(\eta) = \mu^4 \mathcal{P}(z)$ without specifying it in detail. Here $z = \eta/M_{Pl}$ and $\mathcal{P}(z)$ is a function that satisfies all necessary ‘inflationary’ conditions, with the Hubble parameter during inflation $H \sim \mu^2/M_{Pl}$ (for a review, see [21]). The parameter μ is determined by the large scale density perturbations (at scales which reenter the horizon in the matter dominated epoch): $\delta\rho/\rho \simeq O(100)(\mu/M_{Pl})^2$, so that the COBE result $\delta\rho/\rho \simeq 5 \cdot 10^{-6}$ leads to $\mu \sim 10^{15.5}$ GeV. In the context of generic inflation models this in turn implies that the inflaton mass is $m_\eta \sim \mu^2/M_{Pl} \sim 10^{12}$ GeV.

Discussing the couplings of η to the matter fields, we take into the account that in many inflationary scenarios one deals with a large ($\sim M_{Pl}$) amplitude of the inflaton field and moreover, in many cases its VEV is also $\sim M_{Pl}$. Thus in general one cannot neglect higher order terms in η and we formally sum up all them as follows

$$\mu^2 F(z)(\phi^2 + \phi'^2) + \mu^2 \tilde{F}(z)(\phi^2 - \phi'^2) + K(z)(\phi^4 + \phi'^4) + \tilde{K}(z)(\phi^4 - \phi'^4) + A(z)\phi^2 \phi'^2 + \dots \quad (10)$$

where without loss of generality we take the dimensional parameter as μ , and absorb all uncertainties in the unknown factors among which $F(z)$, $K(z)$, $A(z)$ are even functions of z (vanishing at $z = 0$), and $\tilde{F}(z)$, $\tilde{K}(z)$ are odd functions. Lacking understanding of the theory at Planckian energies, we have no apriori information on the shape of these functions. The effects of the possible kinetic-like terms $G(z)[(\partial_\mu \phi)^2 + (\partial_\mu \phi')^2]$ and $\tilde{G}(z)[(\partial_\mu \phi)^2 - (\partial_\mu \phi')^2]$ reduce to redefinition of the wavefunctions of ϕ, ϕ' and for simplicity we do not consider them here. As for the matter fields, we are interested only in their small values, so that the possible higher order terms in ϕ, ϕ' can be neglected.

If η develops a nonzero VEV $\langle \eta \rangle = \eta_0$, then expanding in series of (small) deviations $\tilde{\eta} = \eta - \eta_0$ we see that neither the effective Lagrangian of ϕ and ϕ' , nor their interaction terms with the inflaton field $\tilde{\eta}$ respect P -symmetry anymore:

$$\mathcal{V}(\phi, \phi') = (m^2 \phi^2 + h \phi^4) + (m'^2 \phi'^2 + h' \phi'^4) + a \phi^2 \phi'^2 \quad (11)$$

$$\mathcal{V}(\tilde{\eta}; \phi, \phi') = (f \phi^2 + f' \phi'^2) \mu \tilde{\eta} + (g \phi^2 + g' \phi'^2) \tilde{\eta}^2 + (k \phi^4 + k' \phi'^4) \frac{\tilde{\eta}}{M_{Pl}} + \dots \quad (12)$$

where in general the parameters are all different for the primed and unprimed fields:

$$\begin{aligned} m^2(m'^2) &= m_0^2 + \mu^2(F \pm \tilde{F}), \quad h(h') = h_0 + (K \pm \tilde{K}), \quad a = a_0 + A \\ f(f') &= \frac{\mu}{M_{Pl}}(F_z \pm \tilde{F}_z), \quad g(g') = \frac{\mu^2}{M_{Pl}^2}(F_{zz} \pm \tilde{F}_{zz}), \quad k(k') = K_z \pm \tilde{K}_z \end{aligned} \quad (13)$$

(here the values of the functions F , etc. and their derivatives $F_z = dF/dz$, $F_{zz} = d^2F/dz^2$ etc. are taken at $z_0 = \eta_0/M_{Pl}$). Thus, in principle one can obtain both an asymmetric electroweak breaking and an asymmetric postinflationary reheating. For a certain range of parameters both ϕ and ϕ' would have nonzero and different VEVs, $v' \neq v$. Since P invariance is also broken in the inflaton couplings (12) the widths Γ and Γ' of the decay of $\tilde{\eta}$ into ϕ and ϕ' particles respectively are different. As a result, the reheating temperatures in two worlds $T_R \sim (\Gamma M_{Pl})^{1/2}$ and $T'_R \sim (\Gamma' M_{Pl})^{1/2}$ should also be different.⁵⁾

Let us first consider the case $\eta_0 \sim M_{PL}$. In this case a small size of the VEVs v' and v implies a strong fine tuning. For example, for $v' \simeq 30v$ being about 5 TeV, the values of both F and \tilde{F} should be $\sim 10^{-22}$. Furthermore, F and \tilde{F} should be also fine tuned among each other with the accuracy of 10^{-3} in order to get $v \sim 100$ GeV.

In order to prevent mirror and usual particles from establishing thermal equilibrium with each other, one has to suppress very much the crossing term in (11): $a < O(10)(m'/M_{Pl})^{1/2} \sim 10^{-7}$. The same requirement puts an upper bound on the values of f, f' . Indeed, for energies below the inflaton mass, $m_\eta \sim \mu^2/M_{Pl}$, the first term in eq. (12) mimics the contact term $\delta a \phi^2 \phi'^2$ with $\delta a \sim f f' (\mu/m_\eta)^2$. Thus the above "non-equilibrium" constraint still implies that $f, f' < 10^{-7}$. This limit allows the reheating temperatures T_R and T'_R to be as high as 10^{10} GeV. However, once the functions $F(z)$ and $\tilde{F}(z)$ are small ($\sim 10^{-22}$) at $z = z_0$, it would be unnatural if their derivatives are substantially larger. In other words, without additional fine tunings, these functions and their derivatives should be all very small ($\sim 10^{-22}$) for any z . In this case two particle decays of inflaton into ϕ and ϕ' would lead to unacceptably small T_R and T'_R .

As for K, \tilde{K} and their derivatives K_z, \tilde{K}_z , they are allowed to have values of order unity. Indeed, the coupling constants k, k' (as well as h, h') can be ~ 1 and different from each other. Then 4-body decays of the inflaton lead to different reheating temperatures T_R and T'_R of the typical 'gravitational' size $\sim 0.1(m_\eta^3/M_{Pl})^{1/2} \sim 10^7$ GeV.

However, all these demand a very strong fine tuning between parameters which is not natural. Moreover, in this case everything becomes uncontrollable and arbitrary, P symmetry is actually broken already at the Planck scale and in general it should be violated also in the Yukawa terms due to the big z -induced corrections. In other words, the mirror world becomes a shadow world without any similarity to the visible one.

If η gets VEV at some intermediate scale $\eta_0 \ll M_{Pl}$, for example in the chaotic inflation scenario with simple potential $\mathcal{V}(\eta) = h(\eta^2 - \eta_0^2)^2$ with $h \sim 10^{-15}$ (in the previous notations, this corresponds to $\mathcal{P}(z) = (z^2 - z_0^2)^2$), then the unknown functions in (10) can be expanded in series of z . The interaction terms (10) then can be written as:

$$h_1 \eta^2 (\phi^2 + \phi'^2) + h_2 \eta_0 \eta (\phi^2 - \phi'^2) + h_3 \frac{\eta^2}{M_{Pl}^2} (\phi^4 + \phi'^4) + h_4 \frac{\eta}{M_{Pl}} (\phi^4 - \phi'^4) + \dots \quad (14)$$

One can achieve the (asymmetric) VEVs of ϕ and ϕ' fields in the TeV range and acceptable reheating temperatures T_R and T'_R by choosing a proper range of η_0 and the constants in

⁵⁾ It has been recently emphasized that in some cases parametric resonance may amplify the particle production [22]. While in this case the numerical estimate of the reheating temperature is different, the fact of asymmetric reheating will remain unchanged due to different coupling of the inflaton to the two sectors.

(14), however this still demands unnatural fine tunings. In the next section we present more appealing supersymmetric models for asymmetric inflationary reheating.

4. SUSY models for asymmetric inflationary reheating

In this section we outline a supersymmetric model for chaotic inflation which leads to asymmetric reheating of the normal and the mirror sectors and connects this asymmetry to the different electroweak breakings in the two sectors. We assume the minimal supersymmetric standard model (MSSM) for both sectors denoting the Higgs superfields in the visible and mirror sectors respectively as H_u , H_d and H'_u , H'_d , and introduce also a gauge singlet superfield η . We also assume that P parity is realized as a discrete R symmetry under which $H_u, H_d \rightarrow H'_u, H'_d$, $\eta \rightarrow -\eta$ and $W \rightarrow -W$, and consider the following superpotential:

$$W = \frac{\lambda}{3} \eta^3 + \lambda_1 \eta (H_u H_d + H'_u H'_d) - \lambda_2 \frac{\eta^2}{M_{Pl}} (H_u H_d - H'_u H'_d) \quad (15)$$

We do not include the terms linear in η and bare mass terms of the doublets, assuming that all mass terms arise purely from the soft SUSY breaking scale $m \sim \text{few TeV}$. Then the Higgs potential of the scalar η has the form

$$\mathcal{V}(\eta) = \lambda^2 |\eta|^4 + A m \lambda (\eta^3 + \eta^{*3}) + m^2 |\eta|^2 \quad (16)$$

while the part of the potential involving the Higgs doublets is

$$\begin{aligned} \mathcal{V}(\eta; H, H') &= \left| \lambda_1 - 2\lambda_2 \frac{\eta}{M} \right|^2 |H_u H_d|^2 + \left| \lambda_1 + 2\lambda_2 \frac{\eta}{M} \right|^2 |H'_u H'_d|^2 \\ &+ \left[\left(\lambda_1^2 - 2\lambda_1 \lambda_2 \frac{\eta - \eta^*}{M} + 4\lambda_2^2 \frac{|\eta|^2}{M^2} \right) (H_u H_d)(H'_u H'_d)^* + \text{h.c.} \right] \\ &+ \left| \lambda_1 - \lambda_2 \frac{\eta}{M} \right|^2 |\eta|^2 (|H_u|^2 + |H_d|^2) + \left| \lambda_1 + \lambda_2 \frac{\eta}{M} \right|^2 |\eta|^2 (|H'_u|^2 + |H'_d|^2) \\ &+ \left[\lambda \left(\lambda_1 - 2\lambda_2 \frac{\eta^*}{M} \right) \eta^2 (H_u H_d)^* + \lambda \left(\lambda_1 + 2\lambda_2 \frac{\eta^*}{M} \right) \eta^2 (H'_u H'_d)^* + \text{h.c.} \right] \\ &+ m \left[\left(B\lambda_1 - C\lambda_2 \frac{\eta}{M} \right) \eta H_u H_d + \left(B\lambda_1 + C\lambda_2 \frac{\eta}{M} \right) \eta H'_u H'_d + \text{h.c.} \right] \\ &+ \kappa m^2 (|H_u|^2 + |H_d|^2 + |H'_u|^2 + |H'_d|^2) + \text{D-terms} \end{aligned} \quad (17)$$

where A, B, C, κ are the $O(1)$ constants determined by the SUSY breaking hidden sector.

Clearly, the potential (16) is suitable for chaotic inflation [21] and in order to obtain acceptable density fluctuations we have to assume that $\lambda^2 \sim 10^{-15}$, or $\lambda \sim 10^{-7.5}$. Hence, the scalar part of the η superfield can play a role of the inflaton and its interactions with the Higgs doublets determine the nature of the postinflationary reheating.⁶⁾

⁶⁾ Strictly speaking, the potential (16) occurs in the case of the global supersymmetry. In the minimal supergravity scheme the standard factor $\exp(8\pi|\eta|^2/M_{Pl}^2)$ in the potential would spoil the slow roll condition for $\eta \geq M_{Pl}$. However, one can appeal to the no-scale supergravity scheme with a Kähler potential suggested in [23], in which case the relevant part of the theory works as in a global SUSY case.

If $A > 1$, then the absolute minimum of (16) is achieved for nonzero η : $\eta_0 \sim Am/\lambda \sim 10^{12}$ GeV, which spontaneously breaks P -invariance.⁷⁾ This will immediately induce asymmetry in the electroweak scales. Indeed, substituting $\eta \rightarrow \eta_0$ in the potential (17) we see that the mass terms of the Higgs doublets become different in the visible and mirror sectors. (Notice that μ -terms also are induced, which are asymmetric too.) Taking $\lambda_1 \sim 10^{-7}$ (as well as λ) and $\lambda_2 \sim 1$ (so that $\lambda_2\eta_0/M_{Pl} \sim 10^{-7}$ too), one gets mass terms of H and H' fields in the TeV range but different from each other. Extrapolating this via renormalization group to lower energies we see that the VEVs $\langle H_d \rangle = v_1$ and $\langle H_u \rangle = v_2$ are different from the VEVs of the mirror doublets $\langle H'_d \rangle = v'_1$ and $\langle H'_u \rangle = v'_2$. (In order to obtain the standard electroweak scale $v = (v_1^2 + v_2^2)^{1/2}$ by an order of magnitude smaller than the mirror scale $v' = (v'^2_1 + v'^2_2)^{1/2} \sim \text{few TeV}$, one has to allow ~ 0.1 fine tuning, which seems to be reasonable.) Moreover, in general the ‘up-down’ ratios $\tan \beta = v_2/v_1$ and $\tan \beta' = v'_2/v'_1$ are also different in two sectors. This can alter the content of the mirror baryonic matter discussed in Section 2. In particular, for $\tan \beta' > 2 \tan \beta$ the only stable nucleon in the mirror world would be the mirror neutron.

The second and the third terms in eq. (15) combine to give different decay rates for the inflaton field η into visible and mirror particles, which leads to different reheating temperatures, $T'_R \neq T_R$. Since its mass $m_\eta \sim m \sim \text{few TeV}$, the decay widths into different states are approximately $\sim \lambda^2 m$, which leads to the reheating temperatures T_R, T'_R around a few TeV.

The low reheating temperature excludes many possible mechanisms of baryogenesis but still the electroweak one remains. We have estimated T_R to be of the order of the mirror electroweak scale v' , but much larger than the standard electroweak scale v . Hence, it is likely that after inflationary reheating the mirror universe was already in the phase of the broken electroweak symmetry,⁸⁾ so that the mirror BA might be very small (though the case of a large mirror BA, even larger than the normal one, is by no means excluded). On the other hand, the visible world reheats in unbroken phase. Then in our SUSY model the electroweak phase transition can be easily first order, and the observed BA can be produced due to the (supersymmetric) electroweak baryogenesis mechanism [24]. Moreover, due to supercooling and additional entropy production at the first order phase transition one can suppress the abundance of the mirror particles even if initial reheating temperatures were the same.

In the spirit of this observation, one could in fact replace the last (non-renormalizable) term in (15) by the mass term $\mu(H_u H_d - H'_u H'_d)$ with $\mu \sim m$. This term in combination with the second term in (15) will still cause asymmetry between the VEVs v and v' . However, in this case inflaton couples to H and H' fields in a symmetric way, and the temperature difference between the visible and mirror worlds at the BBN epoch ($T' < T$) can arise purely from the difference in the electroweak scales ($v' > v$), due to the postinflationary ‘miniinflation’ during the (possible) first order phase transition in the visible universe as well as due to the different phase space factors in the inflaton decays

⁷⁾ Notice that there is no problem of domain walls, since the discrete symmetry $\eta \rightarrow -\eta$ is explicitly violated by the trilinear soft term in (16). In addition, an accidental discrete symmetry $\eta \rightarrow \exp(\frac{2\pi i}{3})\eta$ of the potential (16) is also explicitly violated by the third term in the superpotential (15).

⁸⁾ In this case the initial number of species in the mirror world g'_* should be smaller than that in the visible world and the estimates of the Section 2 should be correspondingly changed.

into the usual and mirror particles (since the masses of the latter are closer to the inflaton mass $m_\eta \sim \text{few TeV}$).

Below we present another model which does not have the usual fine tuning problems and can also provide much larger reheating temperatures. Let us consider the supersymmetric $SU(6) \times SU(6)'$ model which could emerge in the context of the $E_8 \times E_8'$ string theory. The SUSY GUT $SU(6)$ [25] has an advantage that it explains the doublet-triplet splitting without fine tuning. The Higgs sector consists of the superfields Σ and $H + \bar{H}$ respectively in adjoint 35 and fundamental $6 + \bar{6}$ representations. The Higgs doublets appear to be light as the pseudo-Goldstone modes of the spontaneously broken accidental global symmetry $SU(6)_\Sigma \times SU(6)_H$ which arises if the crossing terms like $H\Sigma\bar{H}$ are suppressed in the superpotential. At the scale $V_H \sim 10^{17}$ GeV, H, \bar{H} break the local $SU(6)$ symmetry down to $SU(5)$ which is then broken down to $SU(3) \times SU(2) \times U(1)$ by Σ at the scale $V_\Sigma \simeq 10^{16}$ GeV. In this case the Higgs doublets $H_{u,d}$ dominantly come from the doublet fragments of Σ while in H, \bar{H} they are contained with the weight V_Σ/V_H , and the observed hierarchy of fermion masses can be naturally explained in terms of the small parameter $V_\Sigma/V_H \sim 0.1$ [25]. Needless to say that the mirror group $SU(6)'$ is assumed to be exactly the same.

For inflation we take the model similar to the one suggested in [26]. It includes the P -even (S) and P -odd (η) singlet superfields. We assume the following superpotential:

$$W = kS(\eta^2 - \mu^2) + aS\frac{\eta}{M}(\bar{H}H - \bar{H}'H') + bS\frac{\eta^2}{M^2}(\Sigma^2 + \Sigma'^2) + \dots \quad (18)$$

where $\mu \sim 10^{15.5}$ GeV and $k, a, b \sim 10^{-2}$. The vacuum state is $\langle \eta \rangle = \mu$ and $\langle S \rangle = 0$. The tree level potential of inflaton S appears to be flat for large values of the field $S > \mu$. However radiative effects remove the degeneracy and provide necessary “inflationary” profile [26]. The superpotential (18) has an advantage that the slow roll conditions are satisfied for the values of the inflaton field smaller than M_{Pl} ($S \sim 10^{17}$ GeV). Hence the Planck scale corrections are irrelevant and the model can be safely incorporated into the minimal supergravity scheme. The last two terms in (18) combine to give different decay rates of the inflaton into usual and mirror Higgs doublets, so that the reheating temperatures T_R and T'_R in two worlds are different and have the typical magnitude $\sim 10^8$ GeV. It is important to stress that the coupling constants (k, a, b etc.) need not be taken extremely small since the large VEV of η does not induce the large mass terms for the Higgs doublets. The asymmetry of the electroweak scales in two sectors can naturally emerge as a result of the P parity violation in the soft SUSY breaking terms.

5. Conclusions

To summarize, we have discussed cosmological implications of the idea that there is a parallel mirror universe with identical gauge and matter content to the one we inhabit. Consistency with the big bang nucleosynthesis requires that the energy density of the mirror particles should be suppressed with respect to the normal ones. This can be either achieved by a weaker reheating of the mirror world after inflation or by diluting mirror particles by the entropy production in the first order phase transition in the usual world. The latter could be achieved if the usual electroweak scale is below the mirror

one. We have given explicit examples showing how this idea can be realized in realistic models and discussed how this asymmetry may be connected to the different values of the electroweak symmetry breaking scales in the two sectors. At this point our approach drastically differs from the mirror universe models with exact P parity [3, 4]. We have also discussed other cosmological aspects of the mirror universe. In particular, we argue that the mirror baryonic matter is likely to consist only of hydrogen and no heavier nuclei, thus there should be no mirror stars with a thermonuclear active core. There might be cold mirror compact bodies around which could be accessible to microlensing searches. Mirror baryons might form early black holes explaining quasar formation and active galactic nuclei. Mirror baryonic dark matter should be rather different from the usual dark matter because it is dissipative and can possess antibiasing features.

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